

# 750 GeV diphoton resonance in a top and bottom seesaw model

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## Abstract

The top and bottom seesaw model, which extends the top seesaw in order to accomodate a 125 GeV Higgs boson, predicts vector-like top/bottom partners and these partners can be bounded to form several neutral and charged singlet composite scalars by some new strong dynamics. In this letter, we use such a singlet scalar to interpret the 750 GeV diphoton resonance. This singlet scalar is dominantly produced through the gluon fusion process induced by the partners and its diphoton decay is induced by both the partners and the charged singlet scalars. We show that this scenario can readily account for the observed 750 GeV diphoton signal under the current LHC constraints. Further, this scenario predicts some other phenomenology, such as a strong correlation between the decays to  $\gamma\gamma$ ,  $Z\gamma$  and  $ZZ$ , a three-photon signal from the associate production of a singlet scalar and a photon, as well as some signals from the partner cascade decays. These signals may jointly allow for a test of this framework in future 100 TeV hadron collider and ILC experiments.

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## I. INTRODUCTION

The observation of a 125 GeV Higgs boson at the LHC Run-1 [1, 2] is a great triumph of the Standard Model (SM). The current experimental measurements of its production cross sections and decay rates are consistent with the predictions of the SM Higgs boson. However, without a symmetry protection, the SM Higgs mass is quadratically sensitive to the cutoff scale via quantum corrections. This renders the SM rather unnatural and widely motivates new theories beyond the SM. Among many extensions of the SM, the Higgs sector is usually enlarged or modified. So any evidence of non-SM Higgs bosons would indicate the existence of new physics and can be used to elucidate the electroweak symmetry breaking (EWSB) mechanism.

Very recently, the ATLAS and CMS collaborations have reported their first results at 13 TeV LHC and found a resonance-like excess in the diphoton invariant mass spectrum around 750 GeV [3, 4]. The significances of the signals are still only  $3.6\sigma$  and  $2.6\sigma$  in the respective experiments, but if confirmed with more data, this would open the window of new physics at the TeV scale. Several explanations have been proposed for such an excess [5–7]. When interpreting the excess in terms of the production rate of the resonance  $X$ , based on the expected and observed exclusion limits, the CMS and ATLAS experiments at 13 TeV LHC approximately give [6]

$$\sigma_{\gamma\gamma}^{750}(\text{CMS}) = \sigma(pp \rightarrow X) \times Br(X \rightarrow \gamma\gamma) = 5.6_{-2.4}^{+2.4} fb, \quad (1)$$

$$\sigma_{\gamma\gamma}^{750}(\text{ATLAS}) = \sigma(pp \rightarrow X) \times Br(X \rightarrow \gamma\gamma) = 6.0_{-2.0}^{+2.4} fb. \quad (2)$$

Combined with the 8 TeV data [8, 9], the diphoton excess contributing to the combined production rate is given by [6]

$$\sigma_{\gamma\gamma}^{750} = (4.4 \pm 1.1) \text{ fb} . \quad (3)$$

Because of the Landau-Yang theorem [10], the 750 GeV resonance  $X$  can only be a spin-2 or spin-0 particle. However, a graviton-like spin-2 particle with an universal coupling is disfavored by the searches for the  $jj$  [11],  $ZZ$  [12, 13] and  $t\bar{t}$  [14, 15] resonances. Besides, to enhance the diphoton rate, other SM decay modes of the heavy resonance have to be suppressed. So, the most economic way is to construct a theory with a spin-0 SM-singlet scalar  $S$ . Such a singlet naturally has no tree level couplings with the SM particles. While the

large loop couplings  $S_{gg}$  and  $S_{\gamma\gamma}$  can be achieved by introducing new vector-like fermions and/or new charged scalars, which can be found in some composite models and strong dynamics.

In this paper our aim is not to construct a full ultraviolet complete model, but instead work directly with an effective framework inspired by the extension of the top seesaw with the bottom seesaw (namely top and bottom seesaw) [16, 17] to explain the 750 GeV diphoton resonance without conflicting with other LHC data <sup>1</sup>. Because of the heavy mass ( $m_t \sim 175$  GeV), top quark could potentially be associated with the EWSB. The idea of top quark condensation was proposed to explain the EWSB, where a SM Higgs-like  $t\bar{t}$  bound state (called the top-Higgs boson) with a mass  $\sim 2m_t$  is predicted [18]. Obviously, the minimal top condensation model [19] can hardly be consistent with the recent measurements of the Higgs boson at the LHC. To accommodate 125 GeV Higgs boson, some extensions of top quark condensation with seesaw mechanism [16, 17, 20–28] have been widely investigated. Among them, top and bottom seesaw is a feasible way [16, 17]. Such models naturally predict the vector-like top and bottom partners, which can be bounded to form several neutral and charged composite scalars by some new strong dynamics. In our work we use such a neutral singlet scalar (composed of bottom partners) to interpret the 750 GeV resonance. This singlet scalar is dominantly produced through the gluon fusion process induced by the partners and its diphoton decay is induced by both the partners and the charged singlet scalars. Under the current experimental constraints, we find that the 750 GeV diphoton excess can be explained in this top and bottom seesaw scenario.

This paper is organized as follows. In Sec. II, we describe the interactions relevant for the 750 GeV diphoton resonance and discuss the current experimental constraints. In Sec. III we present the numerical results. The conclusion is given in Sec. IV.

## II. THE RELEVANT INTERACTIONS AND CONSTRAINTS

We focus on the relevant interactions for the 750 GeV diphoton resonance within the framework of top and bottom seesaw model [17]. Here we will concern only with the weak isospin singlet sector of the model and decouple it from the electroweak breaking sector

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<sup>1</sup> The original top seesaw model can hardly explain the 750 GeV diphoton excess since the mixing between top quark and top partner usually leads to a sizable branching ratio of the resonance decay to  $t\bar{t}$ .

that is assumed to correctly reproduce the observed Higgs mass. We start from the effective four-fermion interactions (which are assumed to be generated by some strong dynamics at energy scale  $\Lambda$ ) given by

$$\begin{aligned}\mathcal{L}_\Lambda \supseteq & [m_{0\chi}\bar{\chi}_L\chi_R + m_{0\omega}\bar{\omega}_L\omega_R + h.c.] + G_\chi(\bar{\chi}_L\chi_R)(\bar{\chi}_R\chi_L) + G_\omega(\bar{\omega}_L\omega_R)(\bar{\omega}_R\omega_L) \\ & + G_{\chi\omega}(\bar{\omega}_L\chi_R)(\bar{\chi}_R\omega_L) + G_{\omega\chi}(\bar{\chi}_L\omega_R)(\bar{\omega}_R\chi_L),\end{aligned}\quad (4)$$

where  $\chi_{L,R}$  and  $\omega_{L,R}$  are the vector-like top and bottom partners, transforming as singlets under the electroweak  $SU(2)_L$  gauge symmetry. Their SM quantum numbers are given by

$$\chi_L, \chi_R : (3, 1, 2/3), \quad \omega_L, \omega_R : (3, 1, -1/3). \quad (5)$$

At low energy scale  $\mu(< \Lambda)$ , the theory is described in terms of composite fields corresponding to the bounded fermion pairs in Eq.(5). There are six composite scalars relevant for our study, i.e., two neutral singlets  $S_{N_i}$  and four charged singlets  $S_{C_i}^\pm$  ( $i = 1, 2$ )<sup>2</sup>:

$$\begin{aligned}S_{N_1} &\sim \bar{\chi}_L\chi_R, & S_{N_2} &\sim \bar{\omega}_L\omega_R, \\ S_{C_1}^+ &\sim \bar{\omega}_L\chi_R, & S_{C_2}^+ &\sim \bar{\omega}_R\chi_L, & S_{C_1}^- &\sim \bar{\chi}_L\omega_R, & S_{C_2}^- &\sim \bar{\chi}_R\omega_L.\end{aligned}\quad (6)$$

Then, the effective Lagrangian describing the interactions between vector-like quarks and the scalars as well as the self-interactions of the scalars can be written as

$$\begin{aligned}\mathcal{L}_{\mu<\Lambda} \supseteq & y_{N_1}S_{N_1}\bar{\chi}_L\chi_R + y_{N_2}S_{N_2}\bar{\omega}_L\omega_R + y_{C_1}S_{C_1}^+\bar{\omega}_L\chi_R + y_{C_2}S_{C_2}^+\bar{\omega}_R\chi_L \\ & + m_{0\chi}\bar{\chi}_L\chi_R + m_{0\omega}\bar{\omega}_L\omega_R + h.c. + V(S_{N_i}, S_{C_i}^\pm),\end{aligned}\quad (7)$$

where the bare mass terms  $m_{0\chi}$  and  $m_{0\omega}$  are allowed by the SM gauge symmetry. Using large  $N_c$  fermion loop approximation [19], the Yukawa couplings  $y_{N_i}$  at leading order can be estimated as

$$y_{N_i} \simeq \frac{4\pi}{\sqrt{N_c \ln(\Lambda^2/\mu^2)}}. \quad (8)$$

These couplings tend to infinity at the compositeness scale  $\Lambda$  due to the compositeness condition. For example, when  $\Lambda = 10$  TeV,  $\mu = 1$  TeV and  $N_c = 3$ , Yukawa couplings  $y_{N_i} \simeq 3.4$  are predicted. To obtain smaller  $y_{N_i}$ , the cut-off scale  $\Lambda$  should be higher, (in this

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<sup>2</sup> The masses of these composite singlets can be independent of each other since the global symmetry that protects the Higgs boson to be light, is imposed on the electroweak breaking sector [16, 17] and may be broken in the isospin singlet sector.

case, the theory will suffer from the fine tuning, but which is not the focus of this work.), e.g.  $\Lambda = 10^{12}$  TeV,  $\mu = 1$  TeV, then  $y_{N_i} \simeq 1$ . However, it is noted that for  $\Lambda \gg \mu$ , the fermion bubble approximation may not be accurate enough and the full one-loop RG equations are needed to be solved [19]. The potentially large anomalous dimensions can drive large  $y_{N_i}$  values at the compositeness scale down to substantially lower values at low energies [29]. Depending on the details of the full theory, one may in principle end up with hierarchically different Yukawa couplings  $y_{N_i}$  as well. The exact numerical results can be worked out in the full theory, but this is beyond the scope of this paper. After spontaneous symmetry breaking, the vector-like quark masses are  $m_\chi = y_{N_1} \langle S_{N_1} \rangle + m_{0\chi}$  and  $m_\omega = y_{N_2} \langle S_{N_2} \rangle + m_{0\omega}$ . Since the couplings of vector-like quarks to the neutral composite scalars are not proportional to their masses, we can separate the vector-like quark masses from the strength of the interaction  $y_{N_i}$ . This feature can potentially enhance the effective couplings of  $S_{N_i} gg$  and  $S_{N_i} \gamma\gamma$ .

Besides the vector-like quarks, these new charged scalars can contribute to the diphoton decay of  $S_{N_i}$ . The relevant terms of the effective potential  $V(S_{N_i}, S_{C_i}^\pm)$  in Eq.(7) are given by <sup>3</sup>

$$V(S_{N_i}, S_{C_i}^\pm) \supseteq \sum_{i=1}^2 \frac{1}{2} m_{S_{N_i}}^2 S_{N_i}^2 + \sum_{i=1}^2 \frac{1}{2} m_{S_{C_i}^\pm}^2 S_{C_i}^+ S_{C_i}^- + \sum_{i,j=1}^2 \lambda_{C_{ij}} \mu' S_{N_i} S_{C_j}^+ S_{C_j}^-, \quad (9)$$

where  $m_{S_{N_i}}$  and  $m_{S_{C_i}^\pm}$  are the masses of the neutral and charged singlets, respectively.  $\mu'$  is the dimensional parameter and assumed to be 1 TeV. Similarly to Yukawa couplings  $y_{N_i}$ , the trilinear coupling  $\lambda_{C_{ij}}$  can be estimated at leading order through the fermion bubble approximation,

$$\lambda_{C_{ij}} \simeq \frac{32\pi^2}{N_c \ln(\Lambda^2/\mu^2)}, \quad (10)$$

Again, when  $\Lambda \gg \mu$ , the trilinear coupling  $\lambda_{C_{ij}}$  can be reduced. While the neutral scalars develop VEVs, the four-fermion interactions in Eq.(4) resulting in charged scalars are assumed to be sub-critical, in order to avoid spontaneous breaking of  $U(1)_{em}$ .

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<sup>3</sup> In a full theory, the singlets  $S_{N_i}$  may mix with the neutral components of the electroweak doublets. This mixing can be small because the vacuum expectation values  $\langle S_{N_i} \rangle$  can be small and even vanishing. Similarly,  $S_{N_1} H^\dagger H$  and  $S_{N_2} H^\dagger H$  interactions that are respectively induced by  $t - \chi$  and  $b - \omega$  loops can be further suppressed by the large cut-off scale  $\Lambda$  due to the twice transition of  $t - \chi$  and  $b - \omega$ . In this case, the contribution of  $S_{N_i} \rightarrow hh$  channel to the total decay width of  $S_{N_i}$  can be negligibly small. We also require  $m_{S_{C_i}^\pm} > m_{S_{N_i}}/2$  to kinematically forbid the decay channel  $S_{N_i} \rightarrow S_{C_i}^\pm S_{C_i}^\mp$ .

In the top and bottom seesaw framework, the top/bottom quark masses are naturally reduced by the vector-like partners  $\chi/\omega$  via the seesaw in the top/bottom sector. The mass matrices of top and bottom sectors are given by

$$\begin{pmatrix} \bar{t}_L & \bar{\chi}_L \end{pmatrix} \begin{pmatrix} 0 & \mu_1 \\ m_{t\chi} & m_\chi \end{pmatrix} \begin{pmatrix} t_R \\ \chi_R \end{pmatrix}, \quad \begin{pmatrix} \bar{b}_L & \bar{\omega}_L \end{pmatrix} \begin{pmatrix} 0 & \mu_2 \\ m_{b\omega} & m_\omega \end{pmatrix} \begin{pmatrix} b_R \\ \omega_R \end{pmatrix}, \quad (11)$$

where the entries  $\mu_1, \mu_2, m_{t\chi}$  and  $m_{b\omega}$  arise from some strong dynamics. After diagonalizing the above mass matrices with seesaw condition  $m_\chi \gg m_{t\chi}, \mu_1$  and  $m_\omega \gg m_{b\omega}, \mu_2$ , we obtain the physical top/bottom quark masses as

$$m_t^{ph} \approx \frac{\mu_1 m_{t\chi}}{m_\chi}, \quad m_\chi^{ph} \approx m_\chi, \quad (12)$$

$$m_b^{ph} \approx \frac{\mu_2 m_{b\omega}}{m_\omega}, \quad m_\omega^{ph} \approx m_\omega. \quad (13)$$

The mixing angles between top/bottom quarks and their partners are given by

$$\sin^2 \theta_t \approx \frac{m_t^{ph}}{m_\chi}, \quad \sin^2 \theta_b \approx \frac{m_b^{ph}}{m_\omega}. \quad (14)$$

From Eqs.(7) and (14), we can see that the mixing angle  $\sin \theta_t$  cannot be very small, typically  $\mathcal{O}(0.1)$  for a 10 TeV scale vector-like partner due to the large top mass. So, if  $S_{N_1}$  serves as the 750 GeV diphoton resonance, it will have a large branching ratio of  $S_{N_1} \rightarrow t\bar{t}$ , which has been tightly constrained by the null result of the  $t\bar{t}$  resonance search at the LHC Run-1. On the other hand, bottom quark has a small mass and then the mixing angle  $\sin \theta_b$  can be naturally small. Then, if  $S_{N_2}$  is chosen as the 750 GeV diphoton resonance, it can easily satisfy the LHC dijet constraint. So, in our numerical study we require  $m_{S_{N_2}} = 750$  GeV.

### III. NUMERICAL CALCULATIONS AND RESULTS

In Fig.1, we present the Feynman diagrams for the process  $gg \rightarrow S_{N_2}(750 \text{ GeV}) \rightarrow \gamma\gamma$ . The gluon fusion production of  $S_{N_2}$  is induced by the vector-like bottom partner  $\omega$ , while the diphoton decay of  $S_{N_2}$  is induced by both  $\omega$  and the charged scalars  $S_{C_{1,2}}^\pm$ . We calculate the production cross section of  $gg \rightarrow h$  with  $m_h = 750$  GeV at the 13 TeV LHC by using the package HIGLU [30] with CTEQ6.6M PDFs [31]. The renormalization and factorization scales are set as  $\mu_R = \mu_F = m_S/2$ . Then, the cross section of  $gg \rightarrow S_{N_2}$  can be obtained as  $\sigma_{gg \rightarrow S_{N_2}} = (\Gamma_{S_{N_2} \rightarrow gg} / \Gamma_{h_{750} \rightarrow gg}) \cdot \sigma_{gg \rightarrow h_{750}}$ . We also include a  $K$ -factor  $(1 + 67\alpha_s/4\pi)$  [32] in the calculation of the decay width of  $S \rightarrow gg$ .

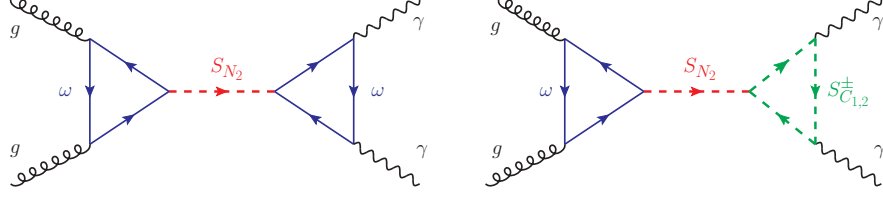


FIG. 1: Feynman diagrams for the process  $gg \rightarrow S_{N_2}(750 \text{ GeV}) \rightarrow \gamma\gamma$ .

The main contributions<sup>4</sup> to the partial width of  $S_{N_2} \rightarrow \gamma\gamma$ ,  $gg$  are from the bottom partner  $\omega$  and charged scalars  $S_{C_{1,2}}^\pm$ , which are given by

$$\Gamma_{\gamma\gamma} \simeq \frac{\alpha^2 m_{S_{N_2}}^3}{256\pi^3} \left| \frac{y_{N_2} \cos^2 \theta_b}{m_\omega} N_c Q_\omega^2 A_{1/2}(\tau_\omega) + \sum_{i=1}^2 \frac{\lambda_{C_{2i}} \mu'}{2m_{S_{C_i}}^2} Q_{S_{C_i}}^2 A_0(\tau_{S_{C_i}}^\pm) \right|^2, \quad (15)$$

$$\Gamma_{gg} \simeq \frac{y_{N_b}^2 \cos^4 \theta_b \alpha_s^2 m_{S_{N_2}}^3}{72\pi^3 m_\omega^2} \left| \frac{3}{4} A_{1/2}(\tau_\omega) \right|^2, \quad (16)$$

where  $N_c = 3$  and  $\tau_\omega = m_{S_{N_2}}^2/4m_\omega^2$ ,  $\tau_{S_{C_i}}^\pm = m_{S_{N_2}}^2/4m_{S_{C_i}}^2$ . The corresponding form factors of the fermion and scalar loops are

$$A_{1/2}(\tau) = 2[\tau + (\tau - 1)f(\tau)]\tau^{-2} \quad (17)$$

$$A_0(\tau) = -[\tau - f(\tau)]\tau^{-2} \quad (18)$$

with

$$f(\tau) = \begin{cases} \arcsin^2 \sqrt{\tau} & \tau \leq 1 \\ -\frac{1}{4} \left[ \log \frac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi \right]^2 & \tau > 1 \end{cases} \quad (19)$$

In our numerical calculations, the input parameters are  $m_{S_{N_i}}$ ,  $m_{S_{C_i}}^\pm$ ,  $m_{\chi, \omega}$ ,  $y_{N_i}$  and  $\lambda_{C_{i1, i2}}$  ( $i = 1, 2$ ). For simplicity, we assume the dimensionless parameters  $\lambda_{C_{11}} = \lambda_{C_{12}} = \lambda_{C_1}$  and  $\lambda_{C_{21}} = \lambda_{C_{22}} = \lambda_{C_2}$  and the charged singlet scalar masses  $m_{S_{C_1}^\pm} = m_{S_{C_2}^\pm} = m_{S_C^\pm}$ . The Yukawa couplings  $y_{N_i}$  and trilinear couplings  $\lambda_{C_{ij}}$  are expected to be large ( $\sim O(1)$ ) but still perturbative ( $\lesssim 4\pi$ ). Since the top partner sector will not contribute to the production rate of the 750 GeV diphoton resonance, we take  $m_{S_{N_1}} = m_\chi = 5 \text{ TeV}$  and assume the corresponding Yukawa coupling as  $y_{S_{N_1}} = 1$  and trilinear couplings  $\lambda_{C_1} = 2$  to avoid the constraints of the LHC search for  $t\bar{t}$  high mass resonance [14, 15] and the electroweak precision observables on top partner sector [24]. It should be noted that the vector-like partner  $\omega$  can provide the radiative corrections to the Peskin-Takeuchi parameter  $T$  [33], and also the corrections to the

<sup>4</sup> The bottom quark contributions is negligible small due to the tiny mixing angle  $\sin \theta_b$ .

$Zb\bar{b}$  vertex induced by the  $b-\omega$  mixing in the bottom seesaw sector. According to Ref. [21], the vector-like bottom partner  $\omega$  should be heavier than 3 TeV to satisfy the requirement of these electroweak precision observables. Then, we scan the relevant parameters in the following ranges

$$375 \text{ GeV} \leq m_{S_{\tilde{C}}^{\pm}} \leq 2 \text{ TeV}, \quad 3 \text{ TeV} \leq m_{\omega} \leq 10 \text{ TeV}, \quad 1 \leq y_{N_2}, \lambda_{C_2} \leq 4\pi. \quad (20)$$

Since there is no resonance observed in the searches for the  $jj$  [11],  $ZZ$  [12, 13] and  $t\bar{t}$  [14, 15], we impose the following constraints in the scan and require our samples to explain the diphoton excess in the  $2\sigma$  range of Eq.(3):

- (1) The CMS search for a dijet resonance [11] at  $\sqrt{s} = 8 \text{ TeV}$  with  $\mathcal{L} = 18.8 \text{ fb}^{-1}$  gives a 95% C.L. upper limit on the production of the RS graviton decaying to  $gg$ ,

$$\sigma(pp \rightarrow X)_{8 \text{ TeV}} \times Br(X \rightarrow gg) < 1.8 \text{ pb} \quad (21)$$

- (2) The ATLAS [12] and CMS [13] searches for a scalar resonance decaying to  $VV$  ( $V = W, Z$ ) at  $\sqrt{s} = 8 \text{ TeV}$  with the full data set, combining all relevant  $Z$  and  $W$  decay channels, give a 95% CL upper limit on the production of the resonance decaying to  $ZZ$ ,

$$\sigma(pp \rightarrow X)_{8 \text{ TeV}} \times Br(X \rightarrow ZZ) < 22 \text{ fb}_{(\text{ATLAS})}, 27 \text{ fb}_{(\text{CMS})} \quad (22)$$

- (3) The ATLAS [8] and CMS [9] searches for a resonance decaying to  $\gamma\gamma$  at  $\sqrt{s} = 8 \text{ TeV}$  give a 95% CL upper limit on the production cross section:

$$\sigma(pp \rightarrow X)_{8 \text{ TeV}} \times Br(X \rightarrow \gamma\gamma) < 2.2 \text{ fb}_{(\text{ATLAS})}, 1.3 \text{ fb}_{(\text{CMS})}. \quad (23)$$

- (4) The run-1 ATLAS resonance search in the  $Z\gamma$  channel give a 95% CL upper limit on the production cross section in the fiducial volume [34]:

$$\sigma(pp \rightarrow X)_{8 \text{ TeV}} \times Br(X \rightarrow Z\gamma) < 4 \text{ fb}_{(\text{ATLAS})}. \quad (24)$$

In Fig.2, we show the scatter plots on the planes of  $m_{S_{\tilde{C}}^{\pm}}$  versus  $\lambda_{C_2}$ , and  $m_{\omega}$  versus  $y_{N_2}$ . The samples are required to satisfy the LHC constraints (1)-(5) and to explain the diphoton excess in the  $2\sigma$  range of Eq.(3). We find that the most stringent constraints on



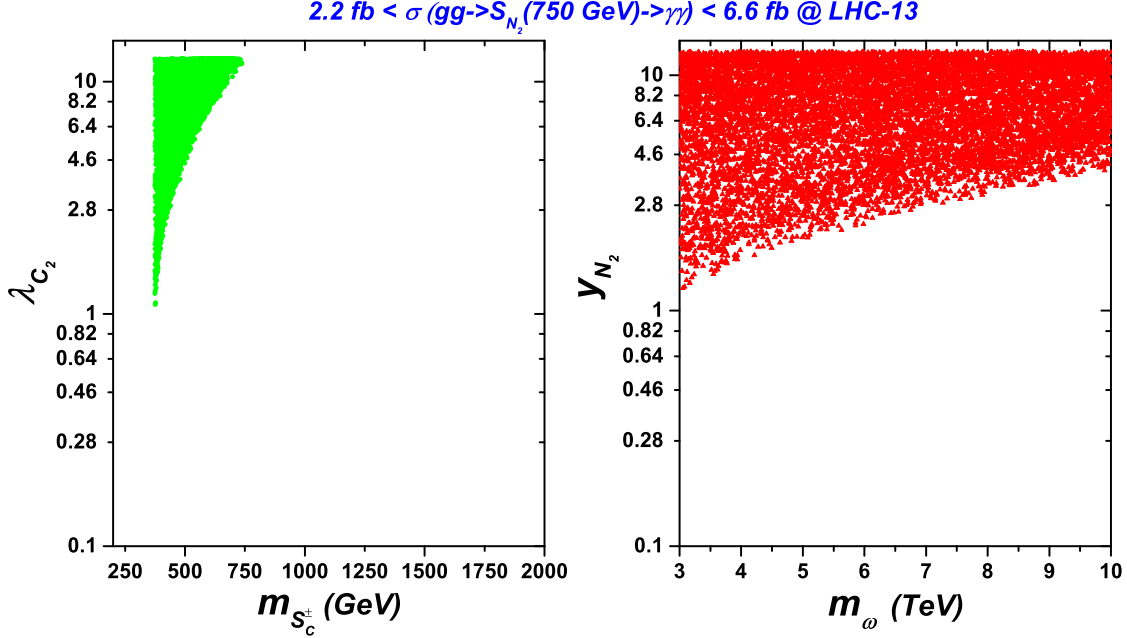


FIG. 2: Scatter plots on the planes of  $m_{S_c^\pm}$  versus  $\lambda_{C_2}$ , and  $m_\omega$  versus  $y_{N_2}$ . The samples satisfy the LHC constraints (1)-(5) and can explain the diphoton excess in the  $2\sigma$  range of Eq.(3).

the parameter space are from the diphoton and dijet measurements at the 8 TeV LHC. The former imposes an upper limit on the cross section of  $gg \rightarrow S_{N_2} \rightarrow \gamma\gamma$  at the 13 TeV LHC, which is about 5.5 fb, while the latter tightly bounds the total decay width  $\Gamma_{tot}$  to be less than about 1 GeV<sup>5</sup>. It is also worth mentioning that the bottom partner  $\omega$  with the Yukawa coupling  $y_{N_2} \gtrsim 1.24$  is mainly responsible for enhancing the cross section of the gluon fusion process  $gg \rightarrow S_{N_2}$ . Only  $\omega$  propagating in the loop of  $S_{N_2} \rightarrow \gamma\gamma$  can hardly explain the

<sup>5</sup> The compatibility of the signal hypothesis depends on the width of the resonance. The CMS data prefers a narrow resonance, while the ATLAS local significance varies from  $3.6\sigma$  for a narrow width to  $3.9\sigma$  for a broader resonance of  $\Gamma_X/m_X \sim 6\%$ . If the ATLAS result persists with more data, the resonance should have a sizable width. In our model, the scalar resonance  $S$  can also provide Majorana type mass for right-handed neutrino  $N$  from the coupling  $\lambda_N S \bar{N}^c N$ . Then we can have the TeV-scale seesaw mechanism in neutrino sector with tiny Yukawa coupling  $y_\nu \bar{L}_L H N_R$ . If the Majorana mass for right handed neutrino is below 375 GeV, the scalar  $S$  can also decay to neutrino pairs and hence the total decay width can be enhanced. Another possibility is to introduce an additional singlet fermion  $N'$ . Then a natural TeV-scale inverse seesaw mechanism can be realized in the neutrino sector. The mass term for  $N'$  can arise from the couplings  $\lambda_{N'} S N' N'$  and can be as light as 100 GeV [35]. So, the scalar  $S$  can also decay to  $S'$  pairs, leading to further enhancement of the decay width. All the additional couplings involving  $S$  can have a dynamical origin. Assuming  $N'$  and  $N$  share other non-critical gauge couplings with  $\chi$ , the four-fermion interactions can generate the desired Yukawa couplings between  $S$  and  $N, N'$  after the condensation of  $\chi$ . In all, the decay width of  $S$  can be easily compatible with the ATLAS data in our model.

diphoton excess since its electric charge is the same as bottom quark,  $Q_\omega = -1/3$ . On the other hand, the lighter charged scalars are, the smaller trilinear couplings are needed. So, from Fig.2 we can see that the light strongly coupled charged scalars with  $m_{S_C^\pm} \lesssim 750$  GeV and  $\lambda_{C_2} \gtrsim 1.1$  are needed to explain the diphoton excess. The heavier mass of  $m_{S_C^\pm}$  is not viable due to the perturbative requirement of the coupling  $\lambda_{C_2}$ .

In our scenario, besides the diphoton excess, other interesting signatures are predicted at the LHC and ILC. For example, due to gauge symmetry, there is a strong correlation between the decay branching ratios of  $S_{N_2} \rightarrow \gamma\gamma$ ,  $S_{N_2} \rightarrow Z\gamma$  and  $S_{N_2} \rightarrow ZZ$ , which is  $1 : 2s_W^2/c_W^2 : s_W^4/c_W^4$ . Moreover, the Drell-Yan process  $q\bar{q}/e^+e^- \rightarrow S_{N_2}(\rightarrow \gamma\gamma)\gamma$  can produce three distinctive hard photons in the final states and may be observed by the future LHC [37] and ILC [38] experiments. Since the bottom partner  $\omega$  is heavier than the 750 GeV resonance  $S_{N_2}$ , we can also have the pair production of  $\omega$ , which leads to  $2b$ -jets and four photons through the cascade decay process  $pp \rightarrow \omega\bar{\omega} \rightarrow S_{N_2}bS_{N_2}\bar{b} \rightarrow 2b + 4\gamma$  at the future 100 TeV hadron collider. Given the current limited significance of diphoton excess, with more data in LHC run-2, both ATLAS and CMS analyses will be able to confirm this excess if it is indeed a signal of new physics beyond the SM. These signatures may be helpful to further test our model at the LHC.

#### IV. CONCLUSIONS

In this work, we interpreted the recent  $\sim 750$  GeV diphoton excess at the 13 TeV LHC in a top and bottom seesaw model. The neutral singlet composite scalar  $S_{N_2}$  (composed of bottom partners) is chosen to play the role of the 750 GeV resonance, which is dominantly produced through gluon fusion process  $gg \rightarrow S_{N_2}$ . Then, the diphoton decay rate of  $S_{N_2}$  can be greatly enhanced by the charged singlet composite scalars. We find that top and bottom seesaw model can account for the observed 750 GeV signal without conflicting with other LHC constraints if the charged scalars have the mass  $m_{S_C^\pm} \lesssim 750$  GeV and the trilinear coupling  $1.1 \lesssim \lambda_{C_2} \lesssim 4\pi$  and the bottom partner has the Yukawa coupling  $1.24 \lesssim y_{N_2} \lesssim 4\pi$ . Besides, this model predicts other signatures, such as the strong correlation between  $S_{N_2} \rightarrow \gamma\gamma$ ,  $S_{N_2} \rightarrow Z\gamma$  and  $S_{N_2} \rightarrow ZZ$  decays, the three-photon signal ( $q\bar{q}/e^+e^- \rightarrow S_{N_2}(\rightarrow \gamma\gamma)\gamma$ ) and the bottom partner cascade decay ( $pp \rightarrow \omega\bar{\omega} \rightarrow S_{N_2}bS_{N_2}\bar{b} \rightarrow 2b + 4\gamma$ ). If the diphoton excess is further confirmed, these signatures may be helpful to test our model in future 100

TeV hadron collider and ILC experiments.

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- [1] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012).
  - [2] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012).
  - [3] CMS Collaboration [CMS Collaboration], CMS-PAS-EXO-15-004.
  - [4] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2015-081.
  - [5] R. Franceschini *et al.*, arXiv:1512.04933 [hep-ph]; S. Di Chiara, L. Marzola and M. Raidal, arXiv:1512.04939 [hep-ph]; K. Harigaya and Y. Nomura, arXiv:1512.04850 [hep-ph]; A. Angelescu, A. Djouadi and G. Moreau, arXiv:1512.04921 [hep-ph]; Y. Nakai, R. Sato and K. Tobioka, arXiv:1512.04924 [hep-ph];
  - [6] D. Buttazzo, A. Greljo and D. Marzocca, arXiv:1512.04929 [hep-ph].
  - [7] T. Higaki, K. S. Jeong, N. Kitajima and F. Takahashi, arXiv:1512.05295 [hep-ph]; S. D. McDermott, P. Meade and H. Ramani, arXiv:1512.05326 [hep-ph]; J. Ellis, S. A. R. Ellis, J. Quevillon, V. Sanz and T. You, arXiv:1512.05327 [hep-ph]; M. Low, A. Tesi and L. T. Wang, arXiv:1512.05328 [hep-ph]; B. Bellazzini, R. Franceschini, F. Sala and J. Serra, arXiv:1512.05330 [hep-ph]; R. S. Gupta, S. Jger, Y. Kats, G. Perez and E. Stamou, arXiv:1512.05332 [hep-ph]; E. Molinaro, F. Sannino and N. Vignaroli, arXiv:1512.05334 [hep-ph].

- ph]; D. Curtin and C. B. Verhaaren, arXiv:1512.05753 [hep-ph]; W. Chao, R. Huo and J. H. Yu, arXiv:1512.05738 [hep-ph]; S. Matsuzaki and K. Yamawaki, arXiv:1512.05564 [hep-ph]; Q. H. Cao, Y. Liu, K. P. Xie, B. Yan and D. M. Zhang, arXiv:1512.05542 [hep-ph]; R. Martinez, F. Ochoa and C. F. Sierra, arXiv:1512.05617 [hep-ph]; B. Dutta, Y. Gao, T. Ghosh, I. Gogoladze and T. Li, arXiv:1512.05439 [hep-ph].
- [8] CMS Collaboration [CMS Collaboration], CMS-PAS-HIG-14-006.
  - [9] G. Aad et al. [ATLAS Collaboration], Phys. Rev. D 92, 032004 (2015) [arXiv:1504.05511 [hep-ex]].
  - [10] L. Landau, Dokl. Akad. Nauk Ser. Fiz. 60 (1948) 207; C. Yang, Phys. Rev. 77 (1950) 242.
  - [11] CMS Collaboration [CMS Collaboration], CMS-PAS-EXO-14-005.
  - [12] G. Aad et al. [ATLAS Collaboration], arXiv:1507.05930 [hep-ex]; arXiv:1509.00389 [hep-ex].
  - [13] V. Khachatryan et al. [CMS Collaboration], JHEP 1510, 144 (2015).
  - [14] G. Aad et al. [ATLAS Collaboration], JHEP 1508, 148 (2015).
  - [15] [CMS Collaboration], CMS-PAS-B2G-12-006.
  - [16] H.-C. Cheng, B. A. Dobrescu, J. Gu, JHEP08(2014)095
  - [17] C. Balazs, T. Li, F. Wang and J. M. Yang, JHEP **1301**, 186 (2013) [arXiv:1208.3767 [hep-ph]].
  - [18] V. Miransky, M. Tanabashi, and K. Yamawaki, Phys. Lett. B221, 177 (1989); V. Miransky, M. Tanabashi, and K. Yamawaki, Mod. Phys. Lett. A4, 1043 (1989); W. Marciano, Phys. Rev. Lett. 62, 2793 (1989); W. J. Marciano, Phys. Rev. D41, 219 (1990);
  - [19] W. A. Bardeen, C. T. Hill, and M. Lindner, Phys. Rev. D41, 1647 (1990).
  - [20] B. A. Dobrescu and C. T. Hill, Phys. Rev. Lett. 81, 2634 (1998); R. S. Chivukula, B. A. Dobrescu, H. Georgi and C. T. Hill, Phys. Rev. D 59, 075003 (1999); B. A. Dobrescu, Phys. Rev. D 63, 015004 (2001).
  - [21] H.-J. He, C. T. Hill and T. M. P. Tait, Phys. Rev. D 65, 055006 (2002).
  - [22] H. S. Fukano and K. Tuominen, Phys. Rev. D **85**, 095025 (2012).
  - [23] H. S. Fukano and K. Tuominen, arXiv:1210.6756 [hep-ph].
  - [24] H. S. Fukano and K. Tuominen, JHEP **1309**, 021 (2013).
  - [25] H. S. Fukano, M. Kurachi, S. Matsuzaki and K. Yamawaki, Phys. Rev. D **90**, no. 5, 055009 (2014).
  - [26] H. S. Fukano and S. Matsuzaki, Phys. Rev. D **90**, no. 1, 015005 (2014).
  - [27] H. S. Fukano, M. Kurachi and S. Matsuzaki, Phys. Rev. D **91**, no. 11, 115005 (2015).

- [28] H.-C. Cheng, J. Gu, JHEP10(2014)002.
- [29] G. Cvetič, Rev. Mod. Phys. **71**, 513 (1999).
- [30] M. Spira, hep-ph/9510347.
- [31] D. Stump *et al.*, JHEP **0310**, 046 (2003) [hep-ph/0303013].
- [32] M. Spira, A. Djouadi, D. Graudenz and P.M. Zerwas, Nucl. Phys. B453 (1995) 17; T. Inami, T. Kubota and Y. Okada, Z. Phys. C18 (1983) 69; A. Djouadi, M. Spira and P.M. Zerwas, Phys. Lett. B264 (1991) 440; M. Spira, A. Djouadi, D. Graudenz and P.M. Zerwas, Phys. Lett. B318 (1993) 347; S. Dawson and R.P. Kauffman, Phys. Rev. D49 (1994) 2298.
- [33] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964 (1990);
- [34] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **738**, 428 (2014).
- [35] P. S. Bhupal Dev, Roberto Franceschini, R. N. Mohapatra, Phys. Rev. D86, 093010 (2012).
- [36] E. Elfgren, hep-ph/0105290.
- [37] R. Brustein and D. H. Oaknin, Phys. Rev. D **62**, 015001 (2000); E. Elfgren, hep-ph/0105290.
- [38] K. Fujii, H. Hano, H. Itoh, N. Okada and T. Yoshioka, Phys. Rev. D **78**, 015008 (2008) doi:10.1103/PhysRevD.78.015008 [arXiv:0802.3943 [hep-ex]].